

An Inflaton Mass Problem in String Inflation from Threshold Corrections to Volume Stabilization

Liam McAllister

*Department of Physics and SLAC
Stanford University, Stanford, CA 94305 USA*

Inflationary models whose vacuum energy arises from a D-term are believed not to suffer from the supergravity eta problem of F-term inflation. That is, D-term models have the desirable property that the inflaton mass can naturally remain much smaller than the Hubble scale. We observe that this advantage is lost in models based on string compactifications whose volume is stabilized by a nonperturbative superpotential: the F-term energy associated with volume stabilization causes the eta problem to reappear. Moreover, any shift symmetries introduced to protect the inflaton mass will typically be lifted by threshold corrections to the volume-stabilizing superpotential. Using threshold corrections computed by Berg, Haack, and Körs, we illustrate this point in the example of the D3-D7 inflationary model, and conclude that inflation is possible, but only for fine-tuned values of the stabilized moduli. More generally, we conclude that inflationary models in stable string compactifications, even D-term models with shift symmetries, will require a certain amount of fine-tuning to avoid this new contribution to the eta problem.

1. Introduction

In any model of slow-roll inflation [1], one needs the inflaton potential $V(\phi)$ to be rather flat, as measured by the slow-roll parameters:

$$\epsilon \equiv \frac{M_{\text{p}}^2}{2} \left(\frac{V'}{V} \right)^2 \quad (1.1)$$

$$\eta \equiv M_{\text{p}}^2 \left(\frac{V''}{V} \right) \quad (1.2)$$

where M_{p} is the four-dimensional reduced Planck mass and primes denote derivatives with respect to the inflaton ϕ . It is convenient to rewrite (1.2) as

$$\eta = \frac{V''}{3H^2} \quad (1.3)$$

so that η measures the inflaton mass in units of the Hubble scale H . Observations require that $\eta \leq 10^{-2}$. A key issue in inflationary model-building is the solution of this constraint.

Inflationary models in supergravity can be divided into F-term models and D-term models according to the source of the supersymmetry-breaking energy which drives inflation. F-term models suffer from what is known as the eta problem, or the inflaton mass problem [2]. The F-term energy

$$V^{\text{F}} \equiv e^K \left(K^{\alpha\bar{\beta}} D_{\alpha} W \overline{D_{\beta} W} - 3|W|^2 \right) \quad (1.4)$$

depends on the inflaton ϕ because ϕ necessarily appears in the Kähler potential. Even if the superpotential depends weakly or not at all on ϕ , the total energy does vary with ϕ . Thus, restoring factors of the Planck mass, we have

$$V_{\text{F}}'' = \frac{K''}{M_{\text{p}}^2} V_{\text{F}} + \dots \quad (1.5)$$

and so a canonically-normalized scalar has $\eta \sim 1$. The only general solution to this problem in F-term models is fine-tuning the contributions in (1.5) to cancel each other to reasonable accuracy, leaving a small net η .

D-term models [3], however, are well-known to be immune to the eta problem, as the Kähler potential does not appear in the D-term energy.¹ This is argued to imply that the

¹ For a discussion of important updates to the D-term inflation scenario, see [4].

inflaton mass need not obey $m_\phi \sim H$, as is generically true in F-term models, but can instead be much smaller. This is a fairly strong argument in favor of D-term inflation.

The goal of this note is to demonstrate that this statement no longer holds in string compactifications whose volume is stabilized by a nonperturbative superpotential: both D-term and F-term models, including shift-symmetric constructions, receive inflaton mass corrections from threshold corrections to the nonperturbative superpotential. We will see that these mass corrections are generically of order the Hubble scale, so that $\eta \sim 1$.

The source of the problem is readily understood. Superpotential stabilization of Kähler moduli proceeds by introducing an F-term potential whose minimum determines² the compactification volume. Just as in the eta problem of F-term inflation, this energy depends on the inflaton through the Kähler potential. Although the inflationary dynamics may be designed to proceed according to a weak interaction, e.g. of widely-separated branes [5,6], the inflaton-dependence of the volume-stabilizing F-term energy typically introduces a *stronger* interaction and renders the total potential too steep for inflation.³

A solution to this problem that has received considerable attention [11,14,15,16] is the introduction of continuous geometric symmetries to protect the inflaton mass. In this approach, one posits the existence of an approximate shift symmetry along the inflationary trajectory.

One purpose of the present paper is to point out that one-loop threshold corrections to the volume-stabilizing nonperturbative superpotentials will typically lift any such shift symmetry and introduce an inflaton mass of order H . Thus, shift symmetries do not suffice to protect the inflaton mass, because quantum corrections will lift these symmetries and change the inflaton potential. Specifically, threshold corrections to the nonperturbative superpotential introduce a dependence of the F-term energy on the various moduli in the system, including both open-string and closed-string fields. The inflaton is usually constructed as one of these moduli, so the F-term potential depends on the inflaton. If this dependence is generic then $\eta \sim 1$. This implies the existence of a rather general eta problem for inflation in nonperturbatively-stabilized string compactifications.

² In some cases, as we will review, the physical volume and the stabilized Kähler modulus are closely related but not identical. For simplicity we will nevertheless refer to this situation as ‘volume stabilization’.

³ This conflict between F-term stabilization and slow-roll inflation was recognized in a concrete form in the brane inflation [5,6] scenario of [7] and has been addressed in e.g. [8,9,10,11,12,13].

Volume stabilization is indispensable for a consistent model, and at present the best-understood methods of volume stabilization use nonperturbative superpotentials, along the lines suggested by Kachru, Kallosh, Linde, and Trivedi (KKLT) [17].⁴ Thus, the presence of an eta problem in the context of nonperturbative volume stabilization is an important aspect of inflation in string theory.

We will be able to observe this effect in detail. Berg, Haack, and Körs (BHK) [8,9] computed the one-loop threshold corrections to the nonperturbative superpotential for the case of type IIB string theory on certain toroidal orientifolds. They observed that the loop corrections introduce a moduli-dependent mass for a mobile D3-brane in this background. (They further showed that this mass correction may be used to fine-tune a brane-antibrane potential to render it flat enough for inflation.) Their result clearly demonstrates, for the case that the inflaton is a D3-brane position and the compactification is a toroidal orientifold, that the inflaton-dependence of the threshold corrections is indeed sufficiently strong to affect inflation.

In §5.2 we will apply the result of BHK to compute the inflaton mass in the D3-D7 inflationary model [19]. A key point is that the D3-D7 model is a D-term model that has been constructed to enjoy a shift symmetry [11], so it might be expected not to be subject to an eta problem. As we will see, even though D-term inflation and shift symmetries do sometimes remove the usual eta problem, neither one suffices to remove the eta problem explored in this paper.

This statement should not be taken as a criticism of the D3-D7 model in particular. We would expect similar results for nearly any model of moving branes in a stabilized string compactification. More generally, the inflaton need not be a brane coordinate; closed string moduli can certainly appear in the threshold corrections, giving a mass to a closed string inflaton. Moreover, although nonperturbative superpotentials play an essential role in our concrete discussion, any F-term moduli-stabilizing energy could in principle lead to the same result.

2. The Eta Problem in Supergravity

In this section we will briefly review the supergravity eta problem and mention how D-term models avoid the problem. In later sections we will argue that this success of D-term models does *not* extend to superpotential-stabilized string compactifications.

⁴ For a very interesting example of perturbative volume stabilization, see [18].

2.1. F-term Inflation and the Eta Problem

In F-term models, inflation proceeds by slowly reducing the F-term energy,

$$V^F \equiv e^K \left(K^{\alpha\bar{\beta}} D_\alpha W \overline{D_\beta W} - 3|W|^2 \right). \quad (2.1)$$

We are interested in computing the slow-roll parameter η (1.2).

Let us work with a canonically normalized inflaton ϕ , which we take to be complex for convenience. Then $\partial_\phi \partial_{\bar{\phi}} K = 1$, so that as a function of ϕ ,

$$V^F(\phi) = V^F(0) \left(1 + \phi \bar{\phi} + \dots \right) \quad (2.2)$$

We may therefore organize the contributions to (1.3) as

$$\eta = 1 + \frac{e^K}{V^F(0)} \partial_\phi \partial_{\bar{\phi}} \left(K^{\alpha\bar{\beta}} D_\alpha W \overline{D_\beta W} - 3|W|^2 \right). \quad (2.3)$$

A successful model requires that the two terms on the right hand side of (2.3) are arranged to cancel to reasonable accuracy, leaving a small net inflaton mass. This sort of fine-tuning is the only *general* solution to the η problem in F-term models.

In particular, if the inflaton does not mix in the Kähler potential with any other fields, so that $K_{,\alpha\bar{\phi}} = 0$ unless $\alpha = \phi$, then the second term in (2.3) depends on the inflaton only through the superpotential, and the necessary fine-tuning must be achieved by adjusting the inflaton-dependence of the superpotential.

2.2. D-term Inflation

D-term models [3,20] are those in which the inflationary trajectory follows a direction which is not D-flat, so that inflation proceeds by slowly reducing a D-term energy. The particular advantage of this approach is that the Kähler potential does not appear in the D-term energy, so the argument of §2.1 does not apply. Thus, the inflaton mass does not receive the corrections of order H that plague F-term models.

At first sight, this conclusion appears surprisingly strong. The mass terms given in §2.1 are merely a concrete example of a general expectation: because the inflationary energy V breaks supersymmetry, we expect soft scalar masses to be induced by gravitational mediation, even if no more direct coupling is present. The resulting masses will be of order $V/M_{\text{p}}^2 = 3H^2$.

Concretely, however, this problematic coupling of the inflaton to the supersymmetry-breaking energy arises from the tree-level Kähler potential for the case of F-term models. D-term inflation sidesteps the problem by providing an inflationary energy which is insensitive to the Kähler potential [21].⁵

We will find that this statement requires careful reexamination in the context of stabilized string compactifications. The reason is that moduli stabilization typically introduces an F-term energy, reviving the problem of §2.1.

3. Nonperturbative Superpotentials and Volume Stabilization

The remainder of our discussion will rely on the details of moduli stabilization, so in this section we will first outline the logic of moduli stabilization and then explain how nonperturbative superpotentials can be used to fix Kähler moduli.

3.1. The Necessity of Volume Stabilization

String compactifications on Calabi-Yau manifolds typically have a large number of massless scalar fields, or moduli. For our purposes the most interesting moduli are the complex structure moduli, the positions of D-branes, and the Kähler parameters, including the overall volume.

Moduli can ruin cosmological models in various ways. They can store energy during inflation and then interfere with nucleosynthesis, or they could have time-dependent vevs at the present epoch, leading to changes in various physical constants. Finally, the presence of these light, gravitationally-coupled fields would typically lead to unobserved fifth-force interactions. Cosmological models which aim to be successful in detail should somehow remove most or all of these light fields.

One modulus in particular presents a grave problem. The overall compactification volume does not have a flat potential, but is in fact unstable: it has a runaway direction toward decompactification. The reason is that the various sources of inflationary energy in string theory will necessarily appear, in the four-dimensional (Einstein-frame) description, multiplied by inverse powers of the volume:

$$V_{4d} = \frac{C}{\rho^\alpha}. \tag{3.1}$$

⁵ S. Thomas has emphasized that Planck-suppressed couplings of the inflaton in the Kähler potential can sometimes produce an inflaton mass even in D-term models [22,23].

Here V_{4d} is the inflationary potential, ρ is the volume modulus (taken to be real), C is a volume-independent factor, and α is positive. This result is easily obtained by dimensional reduction of ten-dimensional sources of energy, such as branes, strings, and fluxes.

If the volume were held fixed by hand, then a mild inflaton-dependence in C could lead to an inflating model. However, in reality we expect that a fast roll in the ρ direction, toward decompactification, will remove the possibility of slow roll in the ϕ direction.

It is therefore absolutely essential to introduce some form of volume-stabilizing potential $U(\rho)$, so that

$$V = \frac{C}{\rho^\alpha} + U(\rho) \quad (3.2)$$

has a minimum at a finite value of ρ .

The proposal of KKLT, which we will now review, is that a nonperturbative superpotential could lead to the necessary volume-dependence.

3.2. Nonperturbative Superpotentials and Volume Stabilization

Let us work in the concrete and well-studied example of the type IIB string on a six-dimensional orientifold, which we view as a limit of a compactification of F-theory on a fourfold. For simplicity we assume that the threefold has exactly one Kähler modulus, ρ . Three-form fluxes H_3, F_3 in the internal space lead to a superpotential [24]

$$W_0 = \int_{\text{CY}} (F_3 - \tau H_3) \wedge \Omega \quad (3.3)$$

which depends on the complex structure moduli $\chi_i, i = 0, \dots, h^{2,1}$ and the dilaton τ .

An additional contribution $W(\rho)$ to the superpotential would allow simultaneous solution of

$$D_\rho W = D_\tau W = D_{\chi_i} W = 0. \quad (3.4)$$

In this supersymmetric solution the dilaton, the complex structure moduli, and the volume are stabilized. (For more details on the stabilization of the complex structure moduli and the dilaton in this scenario, see e.g. [17,25,26].)

KKLT proposed that a nonperturbative superpotential $W_{\text{np}}(\rho)$ from either of two sources could provide the necessary effect:

(1) Euclidean D3-branes wrapping a divisor in the Calabi-Yau [27].

(2) Gaugino condensation on a stack of $N > 1$ D7-branes wrapping a divisor in the Calabi-Yau, and filling spacetime.

In either case, the resulting superpotential takes the form

$$W_{\text{np}} = \Sigma(\zeta, \phi) e^{-a\rho}. \quad (3.5)$$

In this formula a is a numerical constant and Σ is a holomorphic function of the various moduli ζ (such as the complex structure moduli χ_i and the positions of any D-branes) and of the inflaton ϕ .

In the absence of background flux, such a superpotential is possible only when the divisor D satisfies a rather stringent topological condition: the arithmetic genus $\chi(D, \mathcal{O}_D)$ of the divisor must obey $\chi = 1$ [27].

As explained in [28], the effect of fluxes is to permit gaugino condensation to occur somewhat more generally, so that divisors with $\chi > 1$ can contribute to the superpotential. There are reasons to believe that the same conclusion applies to the Euclidean D3-brane superpotential [29,30].

A special feature of the gaugino condensate superpotential is that $a = 4\pi^2/N$ for the condensate of a pure $SU(N)$ gauge group, whereas $a \sim 1$ for the case of Euclidean D3-branes.⁶

We will now turn our attention to the holomorphic prefactor $\Sigma(\zeta, \phi)$.

3.3. Threshold Corrections to Nonperturbative Superpotentials

Recall that in $\mathcal{N} = 1$ Yang-Mills, the Wilsonian gauge coupling is given by the real part of a holomorphic function f :

$$\frac{1}{g^2} = \text{Re}\left(f(\zeta, \phi)\right) \quad (3.6)$$

This holomorphic coupling receives one-loop (and nonperturbative) corrections, but no higher-loop corrections [31,32], so that f is the sum of a tree-level piece and a one-loop correction: $f = f_0 + f_1$.

The one-loop correction f_1 is known as a “threshold correction” because it encodes the effect on the Wilsonian gauge coupling of heavy particles at the threshold, i.e. at the ultraviolet cutoff [33]. This correction is a holomorphic function of the moduli, including, in general, the inflaton.

The gaugino condensate superpotential in pure $SU(N)$ Yang-Mills with ultraviolet cutoff M_{UV} and gauge kinetic function f is given by [31]

⁶ Our conventions for a and ρ differ by a factor of (2π) from those of KKLT: $a_{\text{KKLT}} = 2\pi/N$.

$$W = 16\pi^2 M_{\text{UV}}^3 \exp\left(-\frac{8\pi^2}{N}f\right) \equiv \Sigma(\zeta, \phi)e^{-a\rho}. \quad (3.7)$$

We have absorbed the constants in the exponent into a , we have omitted the dimensionful prefactor, and we have used the fact that dimensional reduction of the $7 + 1$ dimensional theory on the D7-brane relates the tree-level gauge coupling to the volume ρ of the divisor. All further moduli dependence arising from f_1 has been encoded in $\Sigma(\zeta, \phi)$.

In the remainder of the paper we will analyze the physical consequences of the prefactor $\Sigma(\zeta, \phi)$, viewed as a threshold correction to a gaugino condensate superpotential. This means that we are focusing our attention on gaugino condensation instead of Euclidean D3-branes as the source of the superpotential.

The motivation for this choice is that $\Sigma(\zeta, \phi)$ is more readily computed in the gaugino condensate case. For a Euclidean D3-brane superpotential, $\Sigma(\zeta, \phi)$ represents a one-loop determinant of fluctuations around the instanton. In the M-theory description of this effect, this depends on the worldvolume theory of an M5-brane, which is rather subtle [27]. Although explicit results for Σ are unavailable in the Euclidean brane case, we do still expect to find nontrivial inflaton-dependence, leading, as we will see for the gaugino condensate case, to an eta problem.

4. The Eta Problem in String Compactifications

We will now examine the relation between moduli stabilization and the eta problem. In §4.1 we recall a problem which can be thought of as the incarnation of the (usual) supergravity eta problem in a very specific string context. Then, in §4.2 we explain how shift symmetries have been used to address this problem, and we indicate a few important obstacles to the construction of shift-symmetric models.

4.1. Inflaton-Volume Mixing and the Eta Problem

In the context of brane inflation in type IIB string theory, the eta problem takes a novel form [7]. We will examine this now because it presents a concrete setting in which shift symmetries may be used to solve the usual eta problem. Our eventual goal is to understand a new and different eta problem which these symmetries do *not* eliminate, but to achieve this it will be very useful to review the shift symmetry idea in a simpler setting.

D-brane inflation [5] requires mobile, space-filling D-branes, and in a type IIB compactification this is most simply achieved with D3-branes. It will be important for our

considerations that the coordinates of D3-branes (i.e., their center-of-mass position moduli) $\phi_i, i = 1, 2, 3$ appear in the Kähler potential as [34]⁷

$$K = -3 \log \left(\rho + \bar{\rho} - k(\phi_i, \bar{\phi}_i) \right) \quad (4.1)$$

where $k(\phi_i, \bar{\phi}_i)$ is the (unknown) Kähler potential for the Calabi-Yau manifold itself, which is closely related to the D3-brane moduli space. Singling out one direction as the inflaton and denoting it by ϕ , we have $k(\phi, \bar{\phi}) = \phi \bar{\phi} + \dots$, where the expansion is performed around a point in the D3-brane moduli space where the kinetic term is canonical.

This mixing of the brane coordinates with the geometric modulus ρ has important implications. The physical volume r in this setting is no longer simply $Re(\rho)$, but is instead

$$2r = \rho + \bar{\rho} - \phi \bar{\phi}. \quad (4.2)$$

This implies a revision of (3.1), namely

$$V_{4d} = \frac{C}{r^\alpha} = \frac{C}{(\rho - \phi \bar{\phi}/2)^\alpha} \quad (4.3)$$

so that [7]

$$V_{4d}(\phi) = V_{4d}(0) \left(1 + \frac{\alpha}{2r} \phi \bar{\phi} \right) \quad (4.4)$$

This introduces a contribution of order one to η . Because this effect arises from a term in the Kähler potential, it is reasonable to view it as the manifestation, in this specific model, of the usual eta problem. (The new problem we will discuss shortly does not have this property.)

4.2. Solving the Eta Problem with Geometric Shift Symmetries

Shift symmetries [11] are a promising approach to solving the eta problem reviewed in the previous section. The idea is to consider a special compactification which happens to have a particular continuous geometric symmetry.

The proposed symmetry is that the tree-level Kähler potential is independent of one particular (real) field, such as the real part of ϕ . There are strong arguments [15,14] from $\mathcal{N} = 2$ gauged supergravity that this is indeed the case in certain examples, at least before supersymmetry is broken.

⁷ For additional explanation of this point, see [7] and especially [8].

The resulting Kähler potential, for example for D3-branes moving along the torus directions of $K3 \times T^2$, takes the form

$$K = -3 \log \left(\rho + \bar{\rho} - (\phi - \bar{\phi})^2 \right) \quad (4.5)$$

so that $Re(\phi)$ receives no mass from the term analogous to (4.3). This solves the eta problem expressed in (4.4).

Various corrections⁸ will alter this result and lift the shift symmetry of (4.5). In particular, Berg, Haack, and Körs have very clearly demonstrated that threshold corrections to the D7-brane gauge coupling lift the shift symmetry of a certain toroidal orientifold model.

We would like to observe that this conclusion is both generic and problematic, and is in fact a symptom of a new eta problem for string inflation.

Before moving to our main point, we pause to consider some of the obstacles to implementing the shift symmetry argument. (This is an aside because in §5 we will ignore these difficulties and grant the presence of such a symmetry, in the absence of threshold corrections, and then demonstrate that the inclusion of threshold corrections still causes an eta problem.)

The first difficulty is that requiring a geometric shift symmetry places severe constraints on the compactification manifold. It is well-known (cf. [37], p.484) that ordinary Calabi-Yau threefolds, i.e. Calabi-Yau threefolds whose holonomy is $SU(3)$ and not a subgroup, do not have any continuous isometries. Thus, orientifolds of tori and of $K3 \times T^2$ are the only suitable candidates for shift-symmetric models. This implies a tremendous reduction in the number of compactifications available for model-building.

Furthermore, the strategy of guessing general results based on detailed study of toroidal orientifold examples is not always reliable. In particular, even if most such simple examples have continuous symmetries, we know for certain that ordinary Calabi-Yau manifolds do not. Hence, any conclusions about shift symmetries that are inferred from toroidal orientifold examples apply only to that context, and not to the general case. This is one of the reasons that our conclusions are different from those of [16].

Moreover, some important aspects of model-building are actually more difficult in the nominally simplified setting of toroidal orientifolds. Although partial stabilization of

⁸ Perturbative corrections to the Kähler potential will almost certainly lift this symmetry, although we will not address this [35,36].

Kähler moduli has been achieved in this context [38,39], complete stabilization remains challenging. At present it is not clear that known methods will suffice to stabilize all the Kähler moduli in an order-one fraction of toroidal orientifold models. In this regard, Calabi-Yau threefolds with unreduced holonomy can be much more tractable [40]. This is a fairly serious objection to toroidal constructions, given the importance of moduli stabilization for an inflationary model. Even so, it is possible that complete moduli stabilization will eventually be achieved for a toroidal orientifold with properties appropriate for inflation.

5. Threshold Corrections in Nonperturbative Superpotentials Change the Inflaton Mass

We now present the key observation of this paper, which is that threshold corrections induce an entirely new eta problem which D-term and shift-symmetry techniques do *not* solve. That is, we explain how threshold corrections lead to an inflaton mass that is generically of order H , even in the special case that a shift symmetry was present before the inclusion of these corrections.

In §5.1 we discuss the potential sources of an inflaton mass, and in §5.2 we illustrate our considerations with the D3-D7 model [19], in which the problem is particularly clear. In §5.3 we explore potential solutions to this problem.

5.1. General Results

The total potential in a stabilized inflationary model is the sum of several contributions:

$$V = V_F + V_{\text{pos}} + V_{\text{int}}. \quad (5.1)$$

The first contribution, V_F , is the F-term moduli-stabilizing energy. In the KKLT scenario, $V_F = V_{\text{AdS}} < 0$ is also the vacuum energy of a supersymmetric AdS_4 solution. A supersymmetry-breaking effect then adds an energy V_{pos} which ‘uplifts’ the total vacuum energy to a positive value, creating a metastable de Sitter vacuum. The prototypical source of positive energy is an anti-D3-brane [17], though there are various alternatives [41,42].

The final and most model-dependent ingredient is an interaction potential V_{int} designed to produce the dynamics of slow-roll inflation. Simple examples include the weak interactions between a widely-separated brane-antibrane pair [6,7] or between a D3-brane and a D7-brane [19].

The η condition for slow-roll inflation (where primes denote derivative with respect to the canonically-normalized inflaton) is

$$V_F'' + V_{\text{pos}}'' + V_{\text{int}}'' \ll 3H^2 \quad (5.2)$$

By far the simplest case has V_F and V_{pos} independent of ϕ , so that η is determined by V_{int}'' alone. Then, if the interaction potential is reasonably flat, the slow-roll condition can be satisfied. The only remaining challenge is to design an interaction $V_{\text{int}}(\phi)$ that is sufficiently weak.

Of course, this simple case is hard to achieve. Let us now repeat the potential problems:

(1) If V_{int} is an F-term energy then the e^K prefactor leads to an inflaton mass of order H . This is the classic supergravity eta problem [2].

(2) If the inflaton and compactification volume mix, as in (4.2), and the energy is proportional to the volume, as in (4.3), then this produces an eta problem as in (4.4). This was the problem in [7].

(3) If the volume-stabilizing V_F has inflaton dependence, e.g. from threshold corrections, then this leads to yet another eta problem. The inflaton mass depends on the detailed form of these threshold corrections, but is not expected to be parametrically small.

D-term inflationary energy avoids the first problem, as we recalled in §2.2; shift symmetries [11,14] avoid the second problem, as we explained in §4.2; but it appears that some more clever mechanism, or an explicit fine-tuning, will be necessary to overcome the third problem. That is the point of the present note.

5.2. The Example of the D3-D7 Model

It will be worthwhile to illustrate the assertions of the previous section in a specific example. We will focus on the D3-D7 model of [19]. This model is particularly interesting for our purposes because it is a D-term model which can moreover be constructed to take advantage of a shift symmetry, so that the first and second problems of §5.1 are not present. This leaves the inflaton mass from threshold corrections as the final obstacle to a working model.⁹

⁹ Berg, Haack, and Körs have done a careful study [8,9] of the inflaton mass corrections in the brane-antibrane model of [7]. Because the second and third effects listed in §5.1 are *both* present in that example, it is possible to balance these effects against each other and fine-tune away the eta problem. In contrast, our present point is that the third effect, from threshold corrections, is problematic in general, and particularly so in shift-symmetric models.

We will now briefly review the aspects of the D3-D7 model [19]¹⁰ that are relevant for our considerations. The general proposal is that the weak interaction between a mobile D3-brane and a D7-brane whose worldvolume flux \mathcal{F} is not self-dual can give rise to inflation. The D3-brane moves toward the D7-brane and then, at a critical distance, dissolves.

The flux in question is $\mathcal{F} \equiv dA - B$, where A is the gauge potential on the D7-brane worldvolume and B is the pullback of the NS-NS two-form potential. If this flux is not self-dual in the four-dimensional space described by the divisor which the D7-brane wraps, then supersymmetry is broken and there is a force between the D7-brane and the D3-brane [19].

This model can be compactified on $K3 \times T^2/\mathbb{Z}_2$, with the orientifold action explained in [19]. Volume stabilization requires a stack of D7-branes wrapping the $K3$ and sitting at a particular location on the torus, which we take to be the origin. The D7-brane bearing anti-self-dual flux may or may not sit at the same location.

Note that the translational symmetry along the torus may be thought of as the origin of the shift symmetry [11]. Correspondingly, the Kähler potential for this model is given by the shift-symmetric form (4.5)[11,15,14].

The holomorphic gauge coupling on the stack of D7-branes, including the string loop correction, is [9]

$$2f = \rho - \frac{1}{4\pi^2} \log \vartheta_1(\phi, U) + \dots \quad (5.3)$$

where ϑ_1 is a Jacobi theta function, U is the complex structure of the T^2 , ϕ is the inflaton, and the omitted terms are independent of ϕ . Expanding, for convenience, around $\phi = 1/2$, BHK find

$$W_{\text{np}}(1/2 + \phi) = W_{\text{np}}(1/2) \left(1 + \delta(U)\phi^2 \right) \quad (5.4)$$

where

$$\delta(U) = \frac{a}{24} \left(E_2(U) + \vartheta_3(0, U)^4 + \vartheta_4(0, U)^4 \right). \quad (5.5)$$

Here E_2 is the second Eisenstein series, related to derivatives of the ϑ -functions, and a is the numerical constant appearing in (3.5). In these expressions ϕ is dimensionless; the canonically-normalized inflaton, with mass dimension one, is

$$\varphi = M_{\text{p}} \phi \sqrt{\frac{3}{\rho + \bar{\rho}}} \quad (5.6)$$

¹⁰ For a more recent generalization, see [43,44].

We can now compute η for the D3-D7 model on this compactification, by using (5.4) to expand the F-term energy. The result, easily obtained using the SuperCosmology [45] package, is conveniently expressed as

$$\eta = \frac{4}{3} \left| \frac{V_{\text{AdS}}}{V} \right| \left(2 \frac{\delta(U)^2}{a^2} + 3 \frac{\delta(U)}{a} \right) \quad (5.7)$$

where, as in KKLT, V_{AdS} is the vacuum energy at the AdS_4 minimum which is uplifted to create a de Sitter vacuum.

This result is slightly different from the result of [9] for the mass of a $D3 - \overline{D3}$ inflaton. The reason is that the Kähler potential relevant for brane-antibrane inflation is (4.1), but for the present example of D3-D7 inflation the Kähler potential takes the shift-symmetric form (4.5).

Let us now assess whether η (5.7) can satisfy the slow-roll condition $\eta \leq 10^{-2}$. Each of the factors in (5.7), except for $\delta(U)$, is roughly of order one or larger. The ratio $|V_{\text{AdS}}|/V$ cannot be parametrically small, because V_{AdS} determines the height of the potential barrier that prevents decompactification, and the energy density V should not exceed this.¹¹ The constant $1/a$ is likewise not parametrically small; in the concrete example given in KKLT, a was taken to be $2\pi/10$ (where we have included a factor of $(2\pi)^{-1}$ which converts their result to our notation), and more generally, $a = 4\pi^2/N$ for a stack of N coincident D7-branes.

The only factor which might be small is $\delta(U)$. As explained in [9], δ is not automatically small, but there does exist a small range of values of U , the torus complex structure, for which $\delta(U) \ll 1$. A small inflaton mass can therefore be arranged by a choice of fluxes that fixes U in this window. This amounts to an explicit fine-tuning of the inflaton mass.

We conclude that the D3-D7 model requires a modest fine-tuning which can be achieved by a judicious choice of fluxes.¹²

¹¹ This statement is model-dependent; our present discussion assumes moduli stabilization by the method of KKLT [17].

¹² We should again emphasize that corrections to the Kähler potential may introduce further changes in the inflaton mass.

5.3. Discussion

The result of the previous sections accords with the general expectations discussed in §2. An inflaton mass which is much smaller than H does not arise automatically, nor even with the imposition of a shift symmetry; in the end, a fine-tuning at the percent level is necessary to make the model work. In the scheme of inflationary fine-tuning, this is not a serious problem; in particular, it should be contrasted to the functional fine-tuning required for certain models in which $\phi \gg M_{\text{p}}$. Even so, the necessity of fine-tuning in the present case cannot be ignored.

This result should not be interpreted as a stroke against the D3-D7 model (or any other model) in particular. In fact, we would expect almost any complete and fully-realized model to require some fine-tuning of parameters. Omission or simplification of certain physical ingredients, especially moduli stabilization, may obscure the eta problem and make a model appear to work automatically, but sufficient inspection can be expected to reveal one or more problems of detail that require fine-tuning.

It would be extremely interesting to find a solution to this eta problem that does not amount to a fine-tuning of parameters. A slightly modified mechanism of volume stabilization, such as the proposal of [46], does alter the mass formula (5.7), but does not naturally produce a small mass. However, it may be possible to invent a method of volume stabilization which does not affect the inflaton mass. Volume dependence through a D-term energy would be a promising candidate.

Another interesting possibility [21] is that an inflaton charged under a symmetry G can sometimes be excluded from the holomorphic correction term f_1 , so that $\frac{\partial f_1}{\partial \phi} = 0$. However, in simple examples, such as the D3-D7 model, no such symmetry is present. Moreover, D-term inflation requires [3] that ϕ is neutral under the $U(1)$ gauge group G_D whose D-term energy drives inflation, so in particular G cannot coincide with G_D . It is reasonable to expect, however, that discrete symmetries of the appropriate form can sometimes be arranged.

6. Conclusion

We have seen that threshold corrections to volume-stabilizing nonperturbative superpotentials create an eta problem for inflationary models in string theory. These threshold corrections cause the volume-stabilizing F-term energy to depend, generically, on the values of the open-string and compactification moduli. Because the inflaton is expected to

consist of one of these moduli, the threshold correction changes the dependence of the inflationary energy on the inflaton vev, altering the slow-roll parameters and creating an eta problem.

This conclusion applies to models which satisfy several assumptions, which we now repeat for clarity. Our general considerations were limited to models of inflation which can be realized in a string compactification. In any such model it is essential that the instability to decompactification has been removed by moduli stabilization; it is also desirable that all other moduli have also been stabilized. We have explicitly assumed that the volume stabilization arises from a nonperturbative contribution to the superpotential, as in KKLT [17]. (For interesting alternatives, see [18].) We have also assumed that the inflaton is a modulus whose flat direction is slightly lifted by a further supersymmetry-breaking effect. This could correspond, for example, to a brane interaction.

Thus, our result applies to any model of inflation in string theory which uses a compactification stabilized by methods analogous to those of KKLT. Every aspect of the discussion is simplest in the case of D-brane inflation in a type IIB compactification, but the result applies much more broadly. For example, current techniques for moduli stabilization in the heterotic string [47] and in M-theory on G_2 manifolds [48] also use a combination of flux and nonperturbative superpotentials. Any inflationary model¹³ which is elaborated on one of these foundations would be subject to an eta problem from threshold corrections to these superpotentials.

Moreover, although we have seen that the threshold corrections of Berg, Haack and Körs [8] lead to an explicit result for the inflaton mass in the D3-D7 model on $K3 \times T^2$, generic moduli dependence will lead to an eta problem even in more complicated cases. For example, the threshold corrections are not known for generic Calabi-Yau threefolds, so no complete and explicit computation of the slow-roll parameters is possible at present for an inflationary model arising in a compactification on such a space. Progress in this direction appears to be important for inflationary model-building in string compactifications.

It is essential to recognize that although the conclusions of this paper are somewhat general, the actual computation of the inflaton mass is only strictly applicable to a supersymmetric AdS_4 configuration that can be uplifted to produce an inflationary scenario. In particular, the one-loop exactness of threshold corrections in supersymmetric theories permits us to be somewhat precise about the inflaton mass in a supersymmetric vacuum, but,

¹³ For interesting examples in this category, see [49].

as we have emphasized throughout, supersymmetry-breaking effects will typically produce substantial corrections to these mass terms.

Nevertheless, the strategy of understanding the lifting of (inflaton) flat directions in a supersymmetric vacuum is a sensible one.¹⁴ If no suitably flat direction exists in the supersymmetric configuration, it is very hard to believe that the addition of gravitationally-mediated soft terms will remedy this problem. Moreover, it is usually not possible to compute these corrected masses in detail.

Thus, it is usually impossible to prove that a given string model has a small inflaton mass, including all quantum corrections. On the other hand, it is possible to establish that a given model has an eta problem, because if a problem arises from one set of quantum corrections, such as threshold corrections to the gauge coupling, then further quantum corrections will generically not undo this problem. In this paper we have focused on establishing a problem using the one-loop-exact results for the superpotential, with the understanding that additional corrections, e.g. to the Kähler potential, should not conspire to flatten the inflaton potential.

There are several interesting directions for future work. First of all, it is the threshold corrections from closed string moduli that are relevant when the inflaton itself is a closed-string field, for example a geometric modulus [50]. The mass of such a closed string inflaton depends on these corrections, and it would be useful to understand their form.

Furthermore, we have only examined the nonperturbative superpotentials resulting from gaugino condensation, but Euclidean D3-branes are known to play an important role in stabilizing certain classes of Kähler moduli [28,40]. In this context the inflaton dependence of the instanton superpotential arises through a moduli-dependent one-loop determinant $\Sigma(\zeta, \phi)$ of fluctuations around the instanton. It would be extremely interesting, although challenging [27], to compute prefactors of this sort, not only for the considerations of this paper, but for rather general moduli stabilization.

In addition, corrections to the Kähler potential can further adjust the dependence of the total inflationary energy on the inflaton vev. A complete and consistent model requires inclusion of these effects, which have also not yet been calculated.

Looking forward, we can hope that a thorough understanding of the effect of threshold corrections on shift-symmetric brane configurations will guide us to models in which the

¹⁴ This perspective was the one used to expose and address the problem of a brane-antibrane inflaton mass in [7,10,11].

threshold corrections, and all other quantum corrections, are indeed small, so that the shift symmetry is an approximate symmetry of the full quantum theory. If this could be achieved, it would be a significant step toward a controllable model of inflation.

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Appendix A. A Field-Theory Model of the Brane Interaction

In this appendix we will point out a counterintuitive aspect of our conclusion. We will then use a field-theory model to expose the flaw in this intuition, and to further demonstrate that our results are correct.

The inflaton mass term from threshold corrections is the result of an interaction induced by massive strings stretched between the D3-brane and the D7-branes, which we refer to as 3-7 strings. In the field theory description, these 3-7 strings correspond to a massive flavor whose mass m_{37} is controlled by the modulus ϕ . From this perspective, one might expect this massive flavor to decouple when its mass is very large, and to give rise to a negligible interaction in that limit. It is therefore somewhat surprising that the inflaton mass (5.7) does not diminish when $|\phi|$ is large. Should we not expect the BHK result to vanish for widely-separated branes?

To resolve this puzzle, we first note that as soon as m_{37} approaches the mass of a string winding the torus, the D3-D7 interaction induced by the superpotential is correctly described by the full string threshold correction of BHK, and not by its field theory limit. Thus, we can place an upper limit $\Lambda_{UV} < m_W$ on the ultraviolet cutoff of our field-theory description, where m_W denotes the mass of the lightest wound string. In other words, the field theory that provided the decoupling intuition applies only to situations in which the brane separation is much less than the smallest radius of the torus.¹⁵ At greater separations, wound strings can appear in the theory and contribute an additional interaction between the D3-brane and the D7-branes.

We should therefore ask whether decoupling sufficient for slow roll is possible within this limit imposed by the radius of the compact space. To do this, we will examine a simple field theory that models the D3-D7 interaction induced by stretched (but not wound) 3-7 strings. (We will check our model by verifying that it coincides with the small-separation limit of the full BHK result.)

The model is a supersymmetric $SU(N)$ Yang-Mills theory with a single chiral superfield Q whose mass is controlled by a parameter ϕ . Here we will take ϕ to be non-dynamical, and will examine the gaugino condensate superpotential as a function of ϕ .¹⁶

¹⁵ I am grateful to M. Berg and M. Haack for discussions on this point.

¹⁶ For simplicity we are studying the supersymmetric configuration; the supersymmetry-breaking effects used in the model of [19] would generate additional corrections to the inflaton mass, in addition to introducing a tachyon.

The gaugino condensate superpotential below the scale m_{37} , i.e. after integrating out Q , can be matched to the superpotential above this scale. For $N > 2$ the result is simply [51]

$$W_{\text{low}} \propto \Lambda_{\text{high}}^{3-1/N} m_{37}^{1/N} \quad (\text{A.1})$$

with Λ_{high} the dynamically-generated scale of the high-energy theory. Thus, in the low-energy theory, $W = C\phi^{1/N}$ with C independent of ϕ . We can precisely reproduce (A.1) by expanding the superpotential (3.7), including the full string threshold correction (5.3) of BHK, in the limit $\phi \ll 1$, after using the relation $a = 4\pi^2/N$.

Let us now compute η in this model. From the supergravity formula for the F-term energy, we have

$$V = -3e^K C^2 \phi^{2/N} \quad (\text{A.2})$$

so that¹⁷

$$\eta = -\frac{2}{N} \left(1 - \frac{2}{N}\right) \left(\frac{M_{\text{p}}}{\varphi}\right)^2. \quad (\text{A.3})$$

Taking ρ real and using $\varphi = M_{\text{p}}\phi\sqrt{\frac{3}{2\rho}}$, we have

$$\eta \sim -\frac{2}{N} \left(\frac{2\rho}{3\phi^2}\right) \quad (\text{A.4})$$

so that applying $a = 4\pi^2/N$, we finally come to

$$\eta \sim -\frac{a\rho}{3\pi^2\phi^2}. \quad (\text{A.5})$$

However, $a\rho \gg 1$ was a condition for the validity of the nonperturbative superpotential used by KKLT: (3.5) is the leading approximation, analogous to a single-instanton effect, and there will be corrections suppressed by further powers of $e^{-a\rho}$. Furthermore, $|\phi| \leq \frac{1}{2}$ measures the distance from the origin on a unit torus, so $|\phi| \ll \frac{1}{2}$ is necessary in order for the brane separation to be small compared to the size of the torus (and hence for the field theory model to be a good approximation to the true result, which incorporates wound strings.) Thus, there is no controllable parameter regime in which (A.5) is small.

Indeed, even at the extreme boundary of the region of control, $a\rho \sim 1$, $|\phi| \sim \frac{1}{2}$, we have at best $\eta \sim \frac{1}{7}$. If we were to extend the toy model to $\varphi > M_{\text{p}}$ then the interaction would no longer be strong enough to affect slow-roll. However, this is not an allowed range

¹⁷ We have again replaced the dimensionless ϕ with the canonically-normalized φ , cf. (5.6).

in the full model, because of the UV cutoff of the effective field theory, which corresponds to the limit imposed by the radius of the compact space.

We conclude that one cannot arrange suitable decoupling simply by separating the branes; a somewhat more complicated fine-tuning will be necessary to remove the inflaton mass terms under consideration. We have certainly not demonstrated that slow-roll is impossible for D3-D7 systems in the regime in which separations are small compared to the size of the torus. We have simply shown that the interaction captured by threshold corrections produces, on its own, an unsuitably large inflaton mass in this range, so that some fine-tuning against other effects would be needed to make a phenomenologically acceptable model. Thus, one cannot evade the arguments of this paper by separating the D3-brane from the D7-branes and invoking decoupling.

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